

Supersymmetric left-right model and its tests in linear colliders

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Abstract

We investigate phenomenological implications of a supersymmetric left-right model based on $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry testable in the next generation linear colliders. We concentrate in particular on the doubly charged $SU(2)_R$ triplet higgsino $\tilde{\Delta}$, which we find very suitable for experimental search. We estimate its production rate in e^+e^- , e^-e^- , $e^-\gamma$ and $\gamma\gamma$ collisions and consider its subsequent decays. These processes have a clear discovery signature with a very low background from other processes.

1. Introduction

Among the possible extensions of the Standard Model of electroweak interactions perhaps the most appealing one is the left-right symmetric model based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [1]. Apart from its original motivation of providing a dynamical explanation for the parity violation observed in low-energy weak interactions, this model differs from the Standard Model in another important respect: it can explain the observed lightness of neutrinos in a natural way. Neutrino masses are created through the see-saw mechanism [2], according to which there are in each family a light neutrino, much lighter than the charged fermions of the family, and a heavy neutrino. The anomalies measured in the solar [3] and atmospheric [4] neutrino fluxes seem indeed to indicate that neutrinos should have a small but non-vanishing mass. Furthermore, the recent observations of the COBE satellite [5] may indicate that there exists a hot neutrino component in the dark matter of the Universe. The see-saw mechanism can account for all these phenomena, while in the Standard Model neutrinos are massless. In other respects the left-right symmetric model in the low-energy limit is very similar to the Standard Model and is like it in a good agreement with all the laboratory experiments performed so far.

On the technical side, the left-right symmetric model has a naturalness problem similar to that in the Standard Model: the masses of the fundamental Higgs scalars diverge quadratically. To make these divergences cancel one has to fine tune the parameters of the theory to some 28 decimal places. As in the Standard Model, the supersymmetry (susy) can be used to stabilize the scalar masses and cure this hierarchy problem. There are also other arguments in favor of supersymmetry. It may, for example, play a fundamental role in the theory of quantum gravity.

In this paper we shall study some phenomenological aspects of a supersymmetric extension of the left-right symmetric model¹. So far there are no experimental evidence for the right-handed interactions predicted by the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ theory, let alone supersymmetry. Nevertheless, these concepts have so many attractive features that they deserve an experimental and phenomenological

¹Supersymmetric left-right model has been studied also in [6, 7, 8, 9].

scrutiny. The next generation linear electron colliders [10] will provide an excellent environment for such a study as they are planned to operate in the energy range from 0.5 to 2 TeV where new phenomena, such as left-right symmetry and supersymmetry, are expected to manifest themselves.

The left-right symmetric model itself, without supersymmetry, has many interesting predictions, which can be studied in high-energy electron-positron and electron-electron collisions. These have been recently investigated in refs. [11],[12],[13]. In the present paper we will concentrate on the processes, where supersymmetry is involved. We will look for reactions distinctive for the supersymmetric left-right model allowing to distinguish it from the non-susy theory and e.g. the susy version of the Standard Model. (The experimental signatures of the minimal susy Standard Model in linear colliders have been investigated in ref. [14].) In particular we will study the production of the susy partner of the doubly charged Higgs boson, a novel prediction of the model, and the subsequent decays.

The organization of the paper is as follows. In Section 2 we define our susy $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model. We will consider a minimal version of the theory, where the number of Higgs fields is the smallest possible. It turns out that minimal set of scalars consists of two bidoublets transforming as $(\mathbf{2}, \mathbf{2}, 0)$ under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, and two right-handed triplets $(\mathbf{1}, \mathbf{3}, 2)$ and $(\mathbf{1}, \mathbf{3}, -2)$. In Section 3 we investigate the decays of the doubly charged triplet higgsino and the charged sleptons to find experimental signals of the doubly charged triplet higgsino production. In Section 4 we consider various processes in linear colliders where the triplet higgsinos could be produced and calculate their cross sections. A discussion and conclusions are given in Section 5.

2. A Supersymmetric Left-Right Model

Apart from the existence of the superpartners of the ordinary left-right model particles, the most significant difference between the ordinary and the supersymmetric left-right model concerns the Higgs sector. In the non-susy theory the minimal

set of Higgs fields consists of a bidoublet

$$\phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix} = (\mathbf{2}, \mathbf{2}, 0), \quad (1)$$

and a $SU(2)_R$ triplet

$$\Delta = \begin{pmatrix} \frac{1}{\sqrt{2}}\Delta^+ & \Delta^{++} \\ \Delta^0 & -\frac{1}{\sqrt{2}}\Delta^+ \end{pmatrix} = (\mathbf{1}, \mathbf{3}, 2). \quad (2)$$

The bidoublet breaks the $SU(2)_L \times U(1)_Y$ symmetry and thereby gives masses to quarks and charged leptons, as well as to light weak bosons W_1 and Z_1 . The W_1 and Z_1 are, up to a possible small mixing with the right-handed counterparts, the ordinary left-handed weak gauge bosons associated with the symmetry group $SU(2)_L$. The heavy and so far unobserved weak bosons W_2 and Z_2 obtain their masses in the breaking of the $SU(2)_R \times U(1)_{B-L}$ symmetry into $U(1)_Y$, which is caused by a non-vanishing vacuum expectation value of the triplet Higgs field Δ^0 .

If one wanted to stick puritanically in the left-right symmetry of the Lagrangian, one ought to introduce in addition to the bidoublet and the right-handed triplet Higgs fields also a left-handed triplet Higgs multiplet $\Delta_L = (\mathbf{3}, \mathbf{1}, 2)$. This, however, does not have any significant role to play in the dynamics of the theory and it can therefore be left out from the minimal model.

How does the Higgs sector change when one moves to the supersymmetric theory? In supersymmetrization, the cancellation of chiral anomalies among the fermionic partners of the triplet Higgs fields requires that the Higgs triplet Δ is accompanied by another triplet, δ , with opposite $U(1)_{B-L}$ quantum number. Due to the conservation of the $B-L$ symmetry, δ does not couple with leptons and quarks. In the model that we consider, also another bidoublet is added to avoid trivial Kobayashi-Maskawa matrix for quarks. This comes about because supersymmetry forbids a Yukawa coupling where the bidoublet appears as conjugated. The two bidoublets will be denoted by ϕ_u and ϕ_d .

We have chosen the vacuum expectation values for the Higgses, which break the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ into the $U(1)_{em}$, to be as follows

$$\langle \phi_u \rangle = \begin{pmatrix} \kappa_u & 0 \\ 0 & 0 \end{pmatrix}, \quad \langle \phi_d \rangle = \begin{pmatrix} 0 & 0 \\ 0 & \kappa_d \end{pmatrix}, \quad \langle \Delta \rangle = \begin{pmatrix} 0 & 0 \\ v & 0 \end{pmatrix}, \quad \langle \delta \rangle \equiv 0. \quad (3)$$

Here $\kappa_{u,d}$ are of the order of the electroweak scale 10^2 GeV. The vev v of the triplet Higgs has to be much larger in order to have the masses of the new gauge bosons W_2 and Z_2 sufficiently high. With the choice (3) of the vev's the charged gauge bosons do not mix and W_L corresponds to the observed particle. This follows from our choice of giving to one neutral Higgs field in both ϕ_u and ϕ_d a vev equal to zero. This is a simplifying assumption supported by data: the experimental upper limit for the $W_L - W_R$ mixing angle is as small as 0.005 [15].

Whether the set (3) of the vev's as such realizes the minimization of scalar potential may actually be disputable. This question has been discussed in [6, 9]. It was argued in [6] that one needs to take into account the first order radiative corrections, as well as to introduce another pair of Higgs triplets, in order to get at least a local minimum of the scalar potential. In [9] it was noticed that for a region in parameter space also the tree level vacuum is stable, if also one of the remaining electrically neutral scalars, the superpartner of right-handed neutrino ($\tilde{\nu}_R$), is given a non-zero vacuum expectation value. As this matter has little significance for our considerations and results, we will in the following set for simplicity $\langle \tilde{\nu}_R \rangle = 0$.

Given the vev's as in Eq. (3), the masses of the light weak bosons are given by

$$m_{Z_1} = 1/\sqrt{2}\sqrt{(\kappa_u^2 + \kappa_d^2)(g_L^2 + g'^2)}, \quad m_{W_1} = g_L\sqrt{1/2(\kappa_u^2 + \kappa_d^2)}, \quad (4)$$

where $g' = g_R g_V / \sqrt{g_R^2 + g_V^2}$, and the masses of the heavy ones by

$$m_{Z_2} = \sqrt{2}v\sqrt{g_R^2 + g_V^2}; \quad m_{W_2} = g_R\sqrt{1/2(\kappa_u^2 + \kappa_d^2) + v^2}. \quad (5)$$

The masses of the light gauge bosons are well known from the LEP results, $M_{W_1} = 80.22$ GeV and $M_{Z_1} = 91.18$ GeV [16]. The mass constraints for the heavy weak bosons and the bounds on the left-right mixing obtained from the low-energy charged and neutral current data depend on the assumptions one makes. In the case the

gauge coupling constants g_L and g_R of $SU(2)_L$ and $SU(2)_R$, as well as the CKM-matrix and its equivalent in $V + A$ charged current interactions, are kept unrelated, one obtains from the charged current data the bounds [17] $g_L M_{W_2}/g_R \gtrsim 300$ GeV and $g_L \zeta/g_R \lesssim 0.013$, where ζ is the $W_L - W_R$ mixing angle. From neutral current data one can derive the lower bound $M_{Z_2} \gtrsim 400$ GeV for the mass of the new Z -boson and the upper bound of 0.008 for the Z_1, Z_2 mixing angle. CDF experiment at Tevatron has recently obtained the mass limits $M_{W_2} > 520$ GeV and $M_{Z_2} > 310$ GeV [18]. We make the usual assumption that the left and right couplings are equal, $g_R = g_L$. In the numerical evaluations we take also the vacuum expectation values κ_u and κ_d equal. The results we will present are not very sensitive to these assumptions.

At the same time when the right-handed gauge symmetry is broken, the right-handed neutrinos achieve Majorana masses via a lepton number violating $|\Delta L| = 2$ Yukawa coupling $h_{ij} \overline{\nu_{iL}^c} \Delta^0 \nu_{jR}$. The masses are given by a 3×3 matrix m_M . They may be comparable with the heavy weak boson masses M_{W_2} and M_{Z_2} . This leads to the see-saw mechanism which, as mentioned, explains the smallness of the masses of the ordinary left-handed neutrinos. The masses of the light neutrinos are given by

$$m_\nu \simeq -m_D m_M^{-1} m_D^T, \quad (6)$$

where the matrix m_D follows from the Dirac-type Yukawa coupling $f_{ij} \bar{\nu}_{iR} \phi \nu_{jL}$. Very little is known about the Yukawa coupling constants h_{ij} and f_{ij} , but in order to have neutrino mixings they should not be diagonal. Accordingly the triplet Higgs and higgsino couplings are in general flavour changing, which is an obvious advantage concerning the experimental discovery of these particles.

Let us now define our supersymmetric left-right symmetric model. The superpotential is assumed to have the following form:

$$\begin{aligned} W = & h_u^Q \hat{Q}_L^{cT} \hat{\phi}_u \hat{Q}_R + h_d^Q \hat{Q}_L^{cT} \hat{\phi}_d \hat{Q}_R \\ & + h_u^L \hat{L}_L^{cT} \hat{\phi}_u \hat{L}_R + h_d^L \hat{L}_L^{cT} \hat{\phi}_d \hat{L}_R + h_\Delta \hat{L}_R^T i \tau_2 \hat{\Delta} \hat{L}_R \\ & + \mu_1 \text{Tr}(\tau_2 \hat{\phi}_u^T \tau_2 \hat{\phi}_d) + \mu_2 \text{Tr}(\hat{\Delta} \hat{\delta}). \end{aligned} \quad (7)$$

Here $\hat{Q}_{L(R)}$ stands for the doublet of left(right)-handed quark superfields, $\hat{L}_{L(R)}$ stands for the doublet of left(right)-handed lepton superfields, $\hat{\phi}_u$ and $\hat{\phi}_d$ are the two bidoublet Higgs superfields, and $\hat{\Delta}$ and $\hat{\delta}$ the two triplet Higgs superfields. The generation indices of the quark and lepton superfields are not shown. The quantum numbers of the superfields are summarized in Table 1. In our numerical examples we will use for the Yukawa coupling constant $h_\Delta = 0.3$.

In the superpotential (7) the R -parity, $R = (-1)^{3(B-L)+2S}$, is preserved. This ensures that the susy partners with $R = -1$ are produced in pairs and that the lightest supersymmetric particle (LSP) is stable. The parameters μ_i in Eq. (7) are supersymmetric mass parameters. They are usually close to the scale of the soft supersymmetry breaking parameters in order to preserve the naturalness of the theory [19]. In supersymmetric models, which have also a gauge singlet Higgs field, the μ -type terms are generated by giving a vacuum expectation value for the singlet Higgs. We assume here that the parameters $|\mu_i|$ are of the order of the weak scale.

From the superpotential we can calculate the Yukawa interaction terms for the particles. They are given by the general formula [20]

$$\mathcal{L}_{\text{Yukawa}} = -\frac{1}{2}[(\partial^2 W / \partial \varphi_i \partial \varphi_j) \psi_i \psi_j + (\partial^2 W / \partial \varphi_i \partial \varphi_j)^* \bar{\psi}_i \bar{\psi}_j]. \quad (8)$$

In this formula φ_k denote scalar fields and ψ_k fermions of the chiral superfields. For the scalars and the fermions of the gauge superfields there are also non-supersymmetric mass terms, the soft breaking terms [21], given by

$$\mathcal{L}_{\text{soft}} = -\frac{1}{2} \sum_i m_i^2 |\varphi_i|^2 - \frac{1}{2} \sum_\alpha M_\alpha \lambda_\alpha \lambda_\alpha + B\varphi^2 + A\varphi^3 + h.c., \quad (9)$$

where the second sum corresponds to the soft breaking terms for gauginos. The scalar interaction terms, φ^2 and φ^3 , are the quadratic and cubic interaction terms, which are allowed by gauge symmetry for scalars. The scalar masses are found from the scalar potential

$$V = \sum_i \left| \frac{\partial W}{\partial \varphi_i} \right|^2 + \frac{1}{2} \sum_\alpha \left| g_\alpha \sum_{ij} \varphi_i^\dagger T_{ij}^\alpha \varphi_j \right|^2 + V_{\text{soft}}, \quad (10)$$

where V_{soft} is specified by $\mathcal{L}_{\text{soft}}$ in Eq. (9).

In this work we are especially interested in the doubly charged fermions occurring in the Higgs triplet superfields. Their mass matrix is particularly simple, since doubly charged higgsinos do not mix with gauginos. From Eq.(8) one finds the supersymmetric mass terms for the higgsinos,

$$\begin{aligned}\mathcal{L}_{\text{doublet mass}} &= -\mu_1[-\tilde{\phi}_{2u}^0\tilde{\phi}_{1d}^0 + \tilde{\phi}_{1u}^+\tilde{\phi}_{2d}^- + \tilde{\phi}_{2u}^-\tilde{\phi}_{1d}^+ - \tilde{\phi}_{1u}^0\tilde{\phi}_{2d}^0] + h.c. \\ \mathcal{L}_{\text{triplet mass}} &= -\mu_2[\tilde{\Delta}^+\tilde{\delta}^- + \tilde{\Delta}^{++}\tilde{\delta}^{--} + \tilde{\Delta}^0\tilde{\delta}^0] + h.c.\end{aligned}\quad (11)$$

The triplet higgsinos and Higgses have lepton number two. Consequently the final state of the higgsino decay must also have lepton number two in the case of R -parity conservation. The interaction term which includes the strength with which the doubly charged $\tilde{\Delta}$ decays to lepton and slepton is found from Eq. (8) to be

$$\mathcal{L}_{\tilde{\Delta}\tilde{l}} = -2h_{\Delta}\bar{l}^c\tilde{\Delta}\tilde{l}.\quad (12)$$

The other interactions of the doubly charged higgsinos are found from the superfield interaction term $\hat{\varphi}^\dagger e^{2g\hat{\mathcal{V}}}\hat{\varphi}|_{\theta\theta\bar{\theta}\bar{\theta}}$ between the matter superfields $\hat{\varphi}$ and gauge superfield $\hat{\mathcal{V}}$, and they are given by [20]

$$\mathcal{L}_{\text{int,gauge-matter}} = -igT_{ij}^a V_\mu^a \bar{\psi}_i \bar{\sigma}^\mu \psi_j + (ig_a \sqrt{2} T_{ij}^a \varphi_i^* \lambda^a \psi_j + h.c.),\quad (13)$$

where T is the generator of the gauge group.

In unbroken supersymmetry the masses of the leptons, m_ℓ , are equal to the masses of the sleptons. The soft breaking terms provide new mass terms for the scalar particles in the model. The slepton mass matrix is of the general form [22]

$$\begin{pmatrix} L^2 \tilde{m}^2 + m_\ell^2 & A \tilde{m} m_\ell \\ A \tilde{m} m_\ell & R^2 \tilde{m}^2 + m_\ell^2 \end{pmatrix},\quad (14)$$

where L , R , and A are dimensionless constants and \tilde{m} is a mass parameter. These are in principle different for each generation. When compared to the diagonal terms, the off-diagonal mixing terms are small as they are proportional to the lepton mass. The experimental lower limits for the slepton masses are approximately one half of

the LEP beam energy, $m_{\tilde{\ell}} > 43 - 45 \text{ GeV}$ [16]. The squark mass matrices are of the similar form. In unified supersymmetric models the coloured states are heavier than the uncoloured sleptons [23]. We will assume that the squarks are much heavier than the sleptons. This assumption will become important when one considers the decay modes of charginos.

To find the neutralino and chargino masses we need to consider the interaction terms between the superpartners of gauge bosons, the Higgses, and the higgsinos. These are given by

$$\begin{aligned} \mathcal{L}_{\lambda\psi A} = & ig_{B-L}\sqrt{2}v\lambda_{B-L}\tilde{\Delta}^0 + ig_R\sqrt{2}v(\lambda_R^-\tilde{\Delta}^+ - \lambda_R^0\tilde{\Delta}^0) + ig_L\left(\frac{\kappa_u}{\sqrt{2}}\lambda_L^0\tilde{\phi}_{1u}^0 + \kappa_u\lambda_L^+\tilde{\phi}_{2u}^-\right) \\ & -ig_R\left(\frac{\kappa_u}{\sqrt{2}}\tilde{\phi}_{1u}^0\lambda_R^0 + \kappa_u\tilde{\phi}_{1u}^+\lambda_R^-\right) + ig_L\kappa_d\left(\lambda_L^-\tilde{\phi}_{1d}^+ - \frac{1}{\sqrt{2}}\lambda_L^0\tilde{\phi}_{2d}^0\right) \\ & -ig_R\left(\kappa_d\tilde{\phi}_{2d}^-\lambda_R^+ - \frac{\kappa_d}{\sqrt{2}}\tilde{\phi}_{2d}^0\lambda_R^0\right) + h.c. \end{aligned} \quad (15)$$

The soft supersymmetry breaking terms for the gauginos can be written as

$$\mathcal{L}_{\text{soft}} = -1/2\{m_L(\lambda_L^0\lambda_L^0 + 2\lambda_L^+\lambda_L^-) + m_R(\lambda_R^0\lambda_R^0 + 2\lambda_R^+\lambda_R^-) + m_{B-L}\lambda_{B-L}^0\lambda_{B-L}^0\} + h.c. \quad (16)$$

To diagonalize the chargino and neutralino mass matrices we follow the recipe of [20]. We denote $\psi^{+T} = (-i\lambda_L^+, -i\lambda_R^+, \tilde{\phi}_{1u}^+, \tilde{\phi}_{1d}^+, \tilde{\Delta}^+)$ and $\psi^{-T} = (-i\lambda_L^-, -i\lambda_R^-, \tilde{\phi}_{2u}^-, \tilde{\phi}_{2d}^-, \tilde{\delta}^-)$. The chargino mass matrix depends on the following parameters: the soft gaugino masses m_L and m_R , the supersymmetric Higgs masses μ_1 and μ_2 , the vacuum expectation values κ_u , κ_d , and v , and the gauge coupling g_R and g_L . The mass Lagrangian can be written as

$$\mathcal{L}_{\text{chargino mass}} = -\frac{1}{2}(\psi^{+T}\psi^{-T})\begin{pmatrix} 0 & X^T \\ X & 0 \end{pmatrix}\begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} + h.c. \quad (17)$$

For a given set of values for the parameters, one can find numerically the eigenvalues for $X^\dagger X$ and XX^\dagger matrices. The physical charginos $\tilde{\chi}_i^\pm$, $i = 1, \dots, 5$, are found by

multiplying ψ^+ and ψ^- by the corresponding diagonalizing matrices C^\pm :

$$\tilde{\chi}_i^\pm = \sum_j C_{ij}^\pm \psi_j^\pm. \quad (18)$$

Similarly, for neutralinos we denote $\psi^{0T} = (-i\lambda_L^0, -i\lambda_R^0, -i\lambda_{B-L}^0, \tilde{\phi}_{1u}^0, \tilde{\phi}_{2u}^0, \tilde{\phi}_{1d}^0, \tilde{\phi}_{2d}^0, \tilde{\Delta}^0, \tilde{\delta}^0)$.

The neutralino mass matrix depends in addition to the parameters appearing in the chargino case, also on the gaugino mass m_{B-L} and the gauge coupling g_{B-L} . The largeness of the soft gaugino masses determine the nature of the lightest neutralino, but are free parameters. The measured masses of the weak vector bosons give two constraints for the vev's and the gauge couplings. The mass Lagrangian of neutralinos is written as

$$\mathcal{L}_{\text{neutralino mass}} = -\frac{1}{2}\psi^{0T}Y\psi^0 + h.c. \quad (19)$$

One can then find the eigenvalues of $Y^\dagger Y$. Multiplying the ψ^0 by the diagonalizing matrix N gives the physical Majorana neutralinos $\tilde{\chi}_i^0$, $i = 1, \dots, 9$:

$$\tilde{\chi}_i^0 = \sum_j N_{ij}\psi_j^0. \quad (20)$$

For large soft gaugino masses one finds an LSP with a large higgsino component. In the minimal supersymmetric Standard Model this is an unfavoured situation, if one wants to solve the dark matter problem in terms of LSP, since higgsinos annihilate too rapidly [24]. In our case, however, the large higgsino component is the triplet higgsino $\tilde{\delta}^0$, for which the cosmological situation is very different and worth a separate study. The chargino and neutralino masses have also been studied in ref. [8]. In [8] the μ_2 mixing parameter of the triplet Higgses is taken to be zero, which would correspond to massless doubly charged higgsino.

We have calculated numerically the composition of neutralinos and charginos for different values of the parameters. The neutralinos are Majorana particles, whereas the charginos combine together to form Dirac fermions. In Table 2 we give compositions and masses of physical charginos and neutralinos assuming that $m_{W_R} = 500 \text{ GeV}$, the soft supersymmetry breaking parameters are 1 TeV, and $\mu_1 = \mu_2 = 200 \text{ GeV}$.

3. Decay of the triplet higgsino and slepton

Before going to the triplet higgsino production processes we will in this section consider its decay. The allowed decay modes are

$$\begin{aligned}\tilde{\Delta}^{++} \rightarrow & \Delta^{++} \lambda^0, \Delta^+ \lambda^+, \\ & \tilde{\Delta}^+ W_2^+, \\ & \tilde{l}^+ l^+.\end{aligned}\tag{21}$$

In large regions of the parameter space, the kinematically favoured decay mode is $\tilde{\Delta}^{++} \rightarrow \tilde{l}^+ l^+$. This is of course the case only when $m_{\tilde{l}^+} < m_{\tilde{\Delta}^{++}}$ (at least for some lepton flavour), which we will assume in the following. As the mass of the triplet Higgs Δ is of the order of the $SU(2)_R$ breaking scale v [25], the first two decay channels are forbidden energetically in our case of relatively light triplet higgsinos. For the same reason is the channel $\tilde{\Delta}^+ W_2^+$ kinematically disfavoured, since the mass of W_2 is known to be above 0.5 TeV. The decay channel $\tilde{\Delta}^+ W_1^+$ is forbidden in the case of no $W_L - W_R$ mixing. In the following we will assume that $\tilde{\Delta}^{++}$ (and its charge conjugated state $\tilde{\Delta}^{--}$) decay in 100% into the $\tilde{l}l$ final state.

The charged leptons \tilde{l} can decay either to a charged lepton of the same flavour plus a neutralino, to a neutrino plus a chargino, or to a charged gauge boson plus a sneutrino:

$$\tilde{l}^+ \rightarrow l^+ + \tilde{\chi}_i^0, \tag{22}$$

$$\tilde{l}^+ \rightarrow \nu + \tilde{\chi}_i^+, \tag{23}$$

$$\tilde{l}^+ \rightarrow W^+ + \tilde{\nu}. \tag{24}$$

The decay mode (24) is kinematically disfavoured and we do not consider it. As discussed earlier, there are two slepton states of a given flavour, the left-slepton \tilde{l}_L and the right-slepton \tilde{l}_R , which may slightly mix with each other. The decay of the mass eigenstate predominantly the right-slepton into the neutrino channel will in general be kinematically disfavoured or even forbidden because of the heaviness of the right-handed neutrino.

The interaction responsible on the decays (22) and (23) are given by the Lagrangian

$$\begin{aligned}
\mathcal{L}_{\tilde{l}\text{-decay}} &= \frac{1}{2\sqrt{2}} \left[\bar{\tilde{l}}(1 + \gamma_5)(g_L N_{i1} + g_{B-L} N_{i3}) \tilde{\chi}_i^0 \tilde{l}_L \right. \\
&\quad \left. - \bar{\tilde{l}}(1 - \gamma_5)(g_R N_{i2}^* + g_{B-L} N_{i3}^*) \tilde{\chi}_i^0 \tilde{l}_R \right] \\
&\quad - \frac{1}{2} [\bar{\nu}(1 + \gamma_5) g_L C_{i1}^{-*} \tilde{\chi}_i^+ \tilde{l}_L + \bar{\nu}(1 - \gamma_5) g_R C_{i2}^+ \tilde{\chi}_i^+ \tilde{l}_R] + h.c. \\
&\equiv \sum_{i,j} \bar{\tilde{l}}(v_{ij} - a_{ij}\gamma_5) \chi_i^0 \tilde{l}_j + \sum_{i,j} \bar{\tilde{l}}(v'_{ij} - a'_{ij}\gamma_5) \tilde{\chi}_i^+ \tilde{l}_j, \tag{25}
\end{aligned}$$

where θ is the mixing angle between slepton mass eigenstates \tilde{l}_1 and \tilde{l}_2 . The decay width is then given by the formula

$$\Gamma = \frac{1}{4\pi} \sum_{i,j} (|v_{ij}|^2 + |a_{ij}|^2) \frac{(m_{\tilde{l}_i}^2 - m_l^2 - m_{\tilde{\chi}_j}^2)}{m_{\tilde{l}_i}^2} \left[\left(\frac{m_{\tilde{l}_i}^2 + m_l^2 - m_{\tilde{\chi}_j}^2}{2m_{\tilde{\chi}_j}} \right)^2 - m_l^2 \right]^{1/2}. \tag{26}$$

Which of the various decay channels is the dominant one depends on the mass of the decaying slepton. In Fig. 1 the branching ratios of the different channels are plotted as the function of the left-slepton and right-slepton masses (neglecting the slepton mixing). For the left-slepton decay the channel (20) becomes dominant immediately the slepton mass exceeds the mass of the lightest chargino. The chargino has several decay channels, e.g. into a lepton-slepton pair, a W-chargino pair, and a quark-squark pair.

4. Production of the triplet higgsino

The next generation linear electron colliders will, besides the usual e^+e^- reactions, be able to work also in e^-e^- , $e^-\gamma$ and $\gamma\gamma$ modes. The high energy photon beams can be obtained by back-scattering of intensive laser beam on high energy electrons. It turns out that all these collision modes may be useful for investigation of the susy left-right model.

In the following we shall study the following four reactions where the doubly charged higgsinos $\tilde{\Delta}^{\pm\pm}$ are produced:

$$e^+e^- \rightarrow \tilde{\Delta}^{++}\tilde{\Delta}^{--}, \tag{27}$$

$$e^-e^- \rightarrow \tilde{\chi}^0 \tilde{\Delta}^{--}, \quad (28)$$

$$\gamma e^- \rightarrow \tilde{l}^+ \tilde{\Delta}^{--}, \quad (29)$$

$$\gamma\gamma \rightarrow \tilde{\Delta}^{++} \tilde{\Delta}^{--}. \quad (30)$$

We have chosen these reactions for investigation because they all have a clean experimental signature: a few hard leptons and missing energy. Furthermore, they all have very small background from other processes. The fact that $\tilde{\Delta}^{\pm\pm}$ carries two units of electric charge and two units of lepton number and that it does not couple to quarks makes the processes (27) - (30) most suitable and distinctive tests of the susy left-right model.

Reaction $e^+e^- \rightarrow \tilde{\Delta}^{++}\tilde{\Delta}^{--}$

The triplet higgsino pair production in e^+e^- collision occurs through the diagrams presented in Fig. 2, provided of course that these particles are light enough compared with the available collision energy. In contrast with the triplet Higgs fields whose mass is in the TeV scale [25], the mass of the triplet higgsino, $\tilde{\Delta}^{\pm\pm}$, is not strongly constrained. What is known is that since doubly charged fermions have not been seen in present day accelerators, their masses cannot be much below 100 GeV. In the view of our theory, the mass of $\tilde{\Delta}^{--}$ is given by the susy mass parameter μ_2 (see Eq. (11)), which is a free parameter. As we mentioned before, for the reason of naturality its value should not differ too much from the electroweak breaking scale, i.e. $\mu_2 = O(10^2 \text{ GeV})$.

Besides the mass $M_{\tilde{\Delta}^{--}}$, the total cross section of the reaction at a given collision energy depends on the unknown masses of the selectron and the heavier neutral weak boson Z_2 . Of course, the amplitude of the Z_2 mediated reaction is strongly suppressed in comparison with the photon exchange reaction due to the propagator effect and thus the M_{Z_2} dependence of the cross section is quite negligible when the experimental lower limit is taken into account. Note also that the reaction mediated by the lighter weak boson Z_1 is highly suppressed as $\tilde{\Delta}^{--}$ couples to that boson only through the $Z_1 - Z_2$ mixing.

In Fig. 3 the total cross section for the process (27) is presented as a function of the mass of $\tilde{\Delta}^{--}$ for the collision energy of $\sqrt{s} = 1 \text{ TeV}$ and for two values of the

selectron mass, $m_{\tilde{l}} = 200$ GeV and 400 GeV. As can be seen, the cross section is for these parameter values about 0.5 pb and it is quite constant up to the threshold region. To have an estimate for the event rate, one has to multiply the cross section with the branching ratio of the decay channel of the produced higgsinos used for the search. As pointed out earlier, the favoured decay channel may be

$$\tilde{\Delta}^{--} \rightarrow \tilde{l}^{-} l^{-} \rightarrow l^{-} l^{-} \tilde{\chi}^0. \quad (31)$$

Here l can be any of e , μ and τ with practically equal probabilities. The importance of the competing channel with the $\tilde{\Delta}^{+} W^{+}$ final state depends on the mass of the singly charged triplet higgsino $\tilde{\Delta}^{+}$ and the mass of W_R . One may assume that it is close to the mass of the doubly charged higgsino and larger than that of the slepton \tilde{l} , in which case the channel (31) would dominate. In any case the signature of the pair production reaction (27) would be the purely leptonic final state associated with missing energy. The missing energy is carried by neutrinos or neutralinos.

In the Standard Model a final state consisting of four charged leptons and missing energy can result from cascade decays. In the susy left-right model there are, however, some unique final states not possible in the Standard Model, namely those with non-vanishing separate lepton numbers.

Reaction $e^{-}e^{-} \rightarrow \tilde{\chi}^0 \tilde{\Delta}^{--}$

The production of the triplet higgsino $\tilde{\Delta}^{--}$ in electron–electron collision occurs via a selectron exchange in t-channel (see Fig. 4). The cross section is a function of the unknown masses $M_{\tilde{\Delta}^{--}}$ and $m_{\tilde{e}}$. In Fig. 5 the cross section is presented as a function of $M_{\tilde{\Delta}^{--}}$ for two values of the selectron mass, $m_{\tilde{e}} = 200$ GeV and 500 GeV, at the collision energy $\sqrt{s} = 1$ TeV. It is taken into account in this figure that the final state neutralino mass is related to the triplet higgsino mass as they both depend on the parameter μ_2 . The signature of the reaction is a same-sign lepton pair created in the cascade decay (31) of $\tilde{\Delta}^{--}$, associated with the invisible energy carried by neutralinos. As pointed out earlier the two leptons need not be of the same flavour since the $|\Delta L| = 2$ Yukawa couplings are not necessarily diagonal. This may be useful for distinguishing the process from the selectron pair production $e^{-}e^{-} \rightarrow \tilde{e}^{-}\tilde{e}^{-} \rightarrow e^{-}e^{-} + \text{neutralinos}$, which is the leading process for the selectron

production in the susy version of the Standard Model. In the Standard Model the final states $e^-\mu^-$, $e^-\tau^-$ and $\mu^-\tau^-$ are forbidden.

Reaction $\gamma e^- \rightarrow \tilde{l}^+ \tilde{\Delta}^{--}$

The mechanism for producing high-energy photon beams by Compton back-scattering high intensity laser pulses on high energy electron beams was proposed in ref. [26]. The distribution of the energy fraction $y = E_\gamma/E_e$ transferred to the photon in this process is given by [26]

$$P(y) = \frac{1}{N} \left(1 - y + \frac{1}{1 - y} - \frac{4y}{x(1 - y)} + \frac{4y^2}{x^2(1 - y)^2} \right) \quad (32)$$

where

$$x = \frac{4E_e E_{\text{laser}}}{m_e^2} \quad \text{and} \quad 0 \leq y \leq \frac{x}{1 + x}. \quad (33)$$

The factor N is chosen so that $\int dy P(y) = 1$. As discussed in [14], one should tune the laser energy in such a way that $x = 2(\sqrt{2} + 1)$, since for higher x the conversion efficiency will drop considerably due to the possibility of the back-scattered and laser photons to produce e^+e^- pairs. As a result, the hardest photons will have the energy about $0.83E_e$.

There are three Feynman diagrams contributing to the photoproduction reaction (29): electron exchange in s-channel, selectron exchange in t-channel and triplet higgsino exchange in t-channel (see Fig. 6). In Fig. 7 the total cross section is presented as a function of the triplet higgsino mass for the electron-electron center of mass energy $\sqrt{s_{ee}} = 1$ TeV. The cross section is determined by convoluting the photon energy distribution, i.e. $\sigma(s_{ee}) = \int dy P(y) \sigma(s_{e\gamma})$.

The experimental signature of the reaction is three lepton final state associated with missing energy. The positive lepton is any lepton, and the two negative ones can be any combination of the electron, muon and tau, provided the triplet higgsino coupling is not diagonal. A suitable choice of the final state will cut down the Standard Model background coming e.g. from the reaction $e^- \gamma \rightarrow e^- Z^*$. The cross section is above O(100 fb) for a large range of the masses $M_{\tilde{\Delta}^{--}}$ and $m_{\tilde{e}}$, providing hence a good potential for the discovery of $\tilde{\Delta}^{--}$.

Reaction $\gamma\gamma \rightarrow \tilde{\Delta}^{++} \tilde{\Delta}^{--}$

This reaction is an alternative of, but not competitive with, the reaction (27) for producing a doubly charged higgsino pair. Feynman diagram of the process

is presented in Fig. 8. Because the photon energies are not monochromatic but broadly distributed, no sharp threshold will be visible in the production cross section. Moreover, the maximum collision energy will be some 20% less than the e^+e^- energy. On the other hand, the only unknown parameter in the process is the mass $M_{\tilde{\Delta}^{--}}$ as the couplings are completely determined by the known electric charge of the higgsino.

The cross section of the reaction as a function of $M_{\tilde{\Delta}^{--}}$ is given in Fig. 9 for the collision energy $\sqrt{s_{ee}} = 1$ TeV. The experimental signature of the reaction will be of course the same as for the process (27), i.e. four charged leptons associated with missing energy. The cross section is large because of the photon coupling to electric charge.

5. Discussion and conclusions

The left-right symmetric electroweak model based on the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ symmetry has many attractive features. In particular, in the see-saw mechanism it offers a beautiful and very natural explanation for the lightness of the ordinary neutrinos. On the other hand, like in the Standard Model it has a hierarchy problem in the scalar sector, which can be solved by making the theory supersymmetric.

We have investigated in this paper the experimental signatures of the supersymmetric $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model. We have concentrated in the production and decay of the doubly charged $SU(2)_R$ triplet higgsino $\tilde{\Delta}^{++}$. This particle is very suitable for experimental search for many reasons. It is doubly charged, which means that it does not mix with other particles. Consequently its mass is given by a single parameter, the susy Higgs mass μ_2 , which has to be positive, in contrast with μ_1 , which has an undetermined sign. Also the decays of the $\tilde{\Delta}^{++}$ are very limited, since it carries two units of lepton number and it does not couple to quarks. The nonconservation of the separate lepton numbers L_e , L_μ , and L_τ of the $\tilde{\Delta}^{++}$ couplings may also help to distinguish the signal from the background. These separate lepton

number violating couplings can be studied in the slepton pair production, where one of the reaction amplitudes includes $\tilde{\Delta}^{--}$ exchange [27].

We have calculated the production cross sections of $\tilde{\Delta}^{++}$ (and $\tilde{\Delta}^{--}$) in e^+e^- , e^-e^- , $e^-\gamma$ and $\gamma\gamma$ collisions. We pointed out the clear signals of these reactions, which have no substantial background from the Standard Model physics. From the experimental point of view the process $\gamma\gamma \rightarrow \tilde{\Delta}\tilde{\Delta}$ is especially interesting, since it depends only on one parameter, μ_2 , and its cross section is large for $\mu_2 \lesssim 300$ GeV. For larger $\tilde{\Delta}^{++}$ masses the cross sections are still sizable for the other processes. Depending on the situation and the parameters used, the cross sections are in the range $10 \text{ fb} - 1 \text{ pb}$.

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References

- [1] J.C. Pati and A. Salam, Phys. Rev. D 10 (1974) 275;
R.N. Mohapatra and J. C. Pati, Phys. Rev. D11 (1975) 566, 2558;
G. Senjanovic and R. N. Mohapatra, Phys. Rev. D12 (1975) 1502;
R. N. Mohapatra and R. E. Marshak, Phys. Lett. 91B (1980) 222.
- [2] M. Gell-Mann, P. Ramond and R. Slansky, in *Supergravity*, eds. P. van Nieuwenhuizen and D. Z. Freedman (North Holland 1979);
T. Yanagida, in Proceedings of *Workshop on Unified Theory and Baryon Number in the Universe*, eds. O. Sawada and A. Sugamoto (KEK 1979).
- [3] K. Lande et al., in Proc. XXVth Int. Conf. on High Energy Physics, eds. K.K. Phua and Y. Yamaguchi (World Scientific, Singapore, 1991);
K.S. Hirata et al., Phys. Rev. Lett. 66 (1990) 1301;

- K. Nakamura, Nucl. Phys. (Proc. Suppl.) B31 (1993);
A.I. Abazov et al., Phys. Rev. Lett. 67 (1991) 3332;
V.N. Gavrin, in Proc. XXVIth Int. Conf. in High Energy Physics (Dallas 1992),
to appear.
- [4] K.S. Hirata et al., Phys. Lett. B280 (1992) 146;
D. Casper et al., Phys. Rev. Lett. 66 (1993) 2561.
- [5] G.F. Smoot *et al.*, Ast. J. 396 (1992) L1.
- [6] M. Cvetič, J. C. Pati, Phys. Lett. 135B (1984) 57.
- [7] Y. Ahn, Phys. Lett. 149 B (1984) 337;
R. M. Francis, M. Frank, C. S. Kalman, Phys. Rev. D 43 (1991) 2369.
- [8] R. M. Francis, C. S. Kalman and H.N. Saif, Z. Phys. C59 (1993) 655.
- [9] R. Kuchimanchi, R. N. Mohapatra, University of Maryland preprint UMP-PP-93-209 (1993).
- [10] B. Wiik, a talk in the *Workshop on physics and experiments with linear colliders* (Saariselkä, Finland, September 1991), eds. R. Orava, P. Eerola and M. Nordberg (World Scientific, 1992), p. 83.
- [11] J. Maalampi, A. Pietilä and J. Vuori, Nucl. Phys. B 381 (1992) 544, and Phys. Lett. B 297 (1992) 327;
J. Maalampi and A. Pietilä, Z. Phys. C59 (1993) 257;
J. Maalampi, A. Pietilä and M. Raidal, Phys. Rev. D 48 (1993) 4467.
- [12] D. London, G. Belanger and J.N. Ng, Phys. Lett. B188 (1987) 155.
See also, T. Rizzo, Phys. Lett. B116 (1982) 23.
- [13] C.A. Heusch and P. Minkowski, preprint CERN-TH.6606/92 (1993).
- [14] F. Cuypers, G.J. van Oldenborgh and R. Rückl, Nucl. Phys. B409 (1993) 128,
and Nucl. Phys. B409 (1993) 144.

- [15] G. Altarelli, talk given in the International Europhysics Conf. on High Energy Physics, Marseille, July 1993. To appear in proceedings.
- [16] Particle Data Group, *Review of Particle Properties*, Phys. Rev. **D 45** (1992) S1.
- [17] P. Langacker and S. U. Sankar, Phys. Rev. D40 (1989) 1569;
 J. Polak and M. Zralek, Nucl. Phys. B363 (1991) 385;
 G. Beall, M. Bander and A. Soni, Phys. Rev. D48 (1982) 848;
 G. Altarelli et al., Phys. Lett. B261 (1991) 146;
 F. M. Renard and C. Verzegnassi, Phys. Lett. B260 (1991) 225;
 P. Langacker, Phys. Lett. B256 (1991) 277;
 F. del Aguila, J. M. Moreno and M. Quiros, Phys. Lett. B254 (1991) 476.
- [18] F. Abe *et al.*, CDF Collaboration, Phys. Rev. Lett. 68 (1992) 1463.
- [19] See e.g., H.E. Haber, talk at the Workshop on Recent Advances in the Super-world, preprint HARC, SCIPP-93-22 (1993).
- [20] H. E. Haber, G. L. Kane, Phys. Rev. **117** (1985) 75-263.
- [21] L. Girardello, M. T. Grisaru, Nucl. Phys. **B 194** (1982) 65.
- [22] J. Ellis , S. Rudaz, Phys. Lett. **128 B** (1983) 248.
- [23] G.G. Ross and R.G. Roberts, Nucl. Phys. B377 (1992) 571.
- [24] L. Roszkowski, Phys. Lett. B 262 (1991) 59.
- [25] J.F. Gunion, J. Grifols, A. Mendes, B. Kayser, and F. Olness, Phys. Rev. D40 (1989) 1546; N.G. Deshpande, J.F. Gunion, B. Kayser, and F. Olness, Phys. Rev. D44 (1991) 837.
- [26] I.F. Ginzburg, G.L. Kotkin, V.G. Serbo and V.I. Telnov, Nucl. Instr. Meth. 205 (1983) 47.
- [27] Work in progress.

FIGURE CAPTIONS

Figure 1. a) Branching ratios of the left-slepton as a function of the slepton mass for $\mu_2 = 300$ GeV. b) Branching ratios of the left-slepton as a function of the slepton mass for $\mu_2 = 120$ GeV. c) Branching ratios of the right-slepton as a function of the slepton mass for $\mu_2 = 300$ GeV.

Figure 2. Feynman diagrams for the pair production of the doubly charged higgsinos in electron-positron collisions.

Figure 3. Total cross section for the reaction $e^+e^- \rightarrow \tilde{\Delta}^{++}\tilde{\Delta}^{--}$ as a function of the higgsino mass $m_{\tilde{\Delta}^{++}}$ for two values of the selectron mass $m_{\tilde{l}}$ at the collision energy $\sqrt{s} = 1$ TeV.

Figure 4. Feynman diagrams for the production of the doubly charged higgsino in electron-electron collisions.

Figure 5. Total cross section for the reaction $e^-e^- \rightarrow \tilde{\Delta}^{--}\tilde{\chi}^0$ as a function of the higgsino mass $m_{\tilde{\Delta}^{++}}$ for two values of the selectron mass $m_{\tilde{l}}$ at the collision energy $\sqrt{s} = 1$ TeV.

Figure 6. Feynman diagram for the photoproduction of the doubly charged higgsino.

Figure 7. Total cross section for the reaction $\gamma e^- \rightarrow \tilde{\Delta}^{--}\tilde{l}^+$ as a function of the higgsino mass $m_{\tilde{\Delta}^{++}}$ for two values of the selectron mass $m_{\tilde{l}}$ at the electron-electron (positron) collision energy $\sqrt{s_e} = 1$ TeV.

Figure 8. Feynman diagram for the production of the doubly charged higgsinos in photon photon collision.

Figure 9. Total cross section for the reaction $\gamma\gamma \rightarrow \tilde{\Delta}^{--}\tilde{\Delta}^{++}$ as a function of the higgsino mass $m_{\tilde{\Delta}^{++}}$ at the electron-electron collision energy $\sqrt{s_{ee}} = 1$ TeV.

Superfield	Transformation under $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$
Higgs superfields:	
$\hat{\phi}_u = \begin{pmatrix} \hat{\phi}_1^0 & \hat{\phi}_1^+ \\ \hat{\phi}_2^- & \hat{\phi}_2^0 \end{pmatrix}_u$	$(1,2,2,0)$
$\hat{\phi}_d = \begin{pmatrix} \hat{\phi}_1^0 & \hat{\phi}_1^+ \\ \hat{\phi}_2^- & \hat{\phi}_2^0 \end{pmatrix}_d$	$(1,2,2,0)$
$\hat{\Delta} = \begin{pmatrix} \frac{1}{\sqrt{2}}\hat{\Delta}^+ & \hat{\Delta}^{++} \\ \hat{\Delta}^0 & -\frac{1}{\sqrt{2}}\hat{\Delta}^+ \end{pmatrix}$	$(1,1,3,2)$
$\hat{\delta} = \begin{pmatrix} \frac{1}{\sqrt{2}}\hat{\delta}^- & \hat{\delta}^0 \\ \hat{\delta}^{--} & -\frac{1}{\sqrt{2}}\hat{\delta}^- \end{pmatrix}$	$(1,1,3,-2)$
superfields containing quarks and leptons:	
$\hat{Q}_{Li} = \begin{pmatrix} \hat{u}_{Li} \\ \hat{d}_{Li} \end{pmatrix}$	$(3,2,1,1/3)$
$\hat{Q}_{Ri}^c = \begin{pmatrix} \hat{d}_{Ri}^c \\ \hat{u}_{Ri}^c \end{pmatrix}$	$(3^*,1,2,-1/3)$
$\hat{L}_{Li} = \begin{pmatrix} \hat{\nu}_{Li} \\ \hat{e}_{Li} \end{pmatrix}$	$(1,2,1,-1)$
$\hat{L}_{Ri}^c = \begin{pmatrix} \hat{e}_{Ri}^c \\ \hat{\nu}_{Ri}^c \end{pmatrix}$	$(1,1,2,1)$
gauge superfields:	
\hat{G}	$(8,1,1,0)$
\hat{W}_L	$(1,3,1,0)$
\hat{W}_R	$(1,1,3,0)$
\hat{V}	$(1,1,1,0)$

Table 1: The superfields of the supersymmetric left-right model.

chargino χ_i^- (Dirac fermion, 4-component notation)	$m_{\chi_i^-}$ [GeV]
$\begin{pmatrix} -0.99i\lambda_R^- - 0.07\tilde{\phi}_{2d}^- + 0.09\tilde{\delta}^- \\ 0.81i\lambda_R^+ + 0.08\tilde{\phi}_{1u}^+ - 0.58\tilde{\Delta}^+ \end{pmatrix}$	1230
$\begin{pmatrix} -i\lambda_L^- + 0.1\tilde{\phi}_{2u}^- \\ -i\lambda_L^+ + 0.1\tilde{\phi}_{1d}^+ \end{pmatrix}$	1008
$\begin{pmatrix} 0.02i\lambda_R^- - 0.89\tilde{\phi}_{2d}^- - 0.45\tilde{\delta}^- \\ -0.26i\lambda_R^+ - 0.84\tilde{\phi}_{1u}^+ - 0.48\tilde{\Delta}^+ \end{pmatrix}$	212
$\begin{pmatrix} -0.1i\lambda_L^- - \tilde{\phi}_{2u}^- \\ -0.1i\lambda_L^+ - \tilde{\phi}_{1d}^+ \end{pmatrix}$	192
$\begin{pmatrix} 0.11i\lambda_R^- - 0.45\tilde{\phi}_{2d}^- + 0.89\tilde{\delta}^- \\ 0.52i\lambda_R^+ - 0.54\tilde{\phi}_{1u}^+ + 0.66\tilde{\Delta}^+ \end{pmatrix}$	149
neutralino χ_i^0 (Majorana fermion, 2-component notation)	$m_{\chi_i^0}$ [GeV]
$-0.01i\lambda_R^0 + 0.73i\lambda_L^0 - 0.47i\lambda_{B-L}^0 + 0.04\tilde{\phi}_{1u}^0 - 0.01\tilde{\phi}_{2d}^0 + 0.50\tilde{\Delta}^0 + 0.07\tilde{\delta}^0$	1479
$0.87i\lambda_R^0 - 0.25i\lambda_L^0 - 0.40i\lambda_{B-L}^0 - 0.09\tilde{\phi}_{1u}^0 + 0.03\tilde{\phi}_{2d}^0 + 0.01\tilde{\Delta}^0 + 0.002\tilde{\delta}^0$	1009
$0.48i\lambda_R^0 + 0.48i\lambda_L^0 + 0.74i\lambda_{B-L}^0$	1000
$-0.003i\lambda_R^0 + 0.38i\lambda_L^0 - 0.24i\lambda_{B-L}^0 - 0.07\tilde{\phi}_{1u}^0 - 0.02\tilde{\phi}_{2d}^0 - 0.83\tilde{\Delta}^0 + 0.31\tilde{\delta}^0$	529
$0.04i\lambda_R^0 - 0.01i\lambda_L^0 - 0.02i\lambda_{B-L}^0 + 0.71\tilde{\phi}_{1u}^0 + 0.70\tilde{\phi}_{2d}^0 - 0.05\tilde{\Delta}^0 + 0.05\tilde{\delta}^0$	203
$\tilde{\phi}_{1d}^0$	200
$\tilde{\phi}_{2u}^0$	200
$0.08i\lambda_R^0 - 0.02i\lambda_L^0 - 0.03i\lambda_{B-L}^0 + 0.69\tilde{\phi}_{1u}^0 - 0.71\tilde{\phi}_{2d}^0 - 0.06\tilde{\Delta}^0 - 0.06\tilde{\delta}^0$	194
$-0.002i\lambda_R^0 + 0.18i\lambda_L^0 - 0.11i\lambda_{B-L}^0 - 0.03\tilde{\phi}_{1u}^0 + 0.08\tilde{\phi}_{2d}^0 - 0.24\tilde{\Delta}^0 - 0.94\tilde{\delta}^0$	51

Table 2: Physical charginos and neutralinos for $m_{W_R} = 500$ GeV (or $v = 759$ GeV), $\mu_1 = \mu_2 = 200$ GeV and soft gaugino masses of 1 TeV

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